A STRONG ELECTROWEAK SECTOR AT FUTURE LINEAR COLLIDERS *)

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ABSTRACT

An effective lagrangian describing a strong interacting electroweak sector is considered. It contains new vector and axial-vector resonances all degenerate in mass and mixed with W and Z. The model, for large mass of these degenerate gauge bosons, becomes identical to the standard model in the classical limit of infinite Higgs mass. The limits on the parameter space of this model from future e^+e^- colliders are presented.

1 The Model

In view of future projects of e^+e^- linear colliders it is important to study the possible phenomenology at such colliders from a strong electroweak sector [1]. We shall study the effects of the strong electroweak sector af future linear colliders, assuming a low energy effective theory. The effective lagrangian contains vector and axial-vector resonances as the most visible manifestations at low energy of the strong interacting sector [2]. This model is an extension of the BESS model where only new vector resonances are present [3]. It leads to an interesting and appealing phenomenology.

Let us call G the symmetry group of the theory, spontaneously broken. Among the Goldstone bosons, three are absorbed to give mass to W and Z. The vector and axial-vector mesons will transform under the unbroken subgroup H of G. Following the *hidden gauge symmetry* approach [4][5], theories with non linearly realized symmetry $G \to H$ can be linearly realized by enlarging the gauge symmetry G to $G \otimes H' \to H_D = diag(H \otimes H')$. H' is a local gauge group and the vector and axial-vector are the gauge fields associated to H'.

Let us consider such an effective lagrangian parameterization for the electroweak symmetry breaking, using $G = SU(2)_L \otimes SU(2)_R$, $H' = SU(2)_L \otimes SU(2)_R$. The nine Goldstone

bosons resulting from the spontaneous breaking of $G' = G \otimes H'$ to H_D, can be described by three independent $SU(2)$ elements: L, R and M, with the following transformations properties

$$L' = g_L L h_L, \quad R' = g_R R h_R, \quad M' = h_R^\dagger M h_L \tag{1.1}$$

with $g_{L,R} \in SU(2)_{L,R} \subset G$ and $h_{L,R} \in H'$. Moreover we shall require the invariance under the discrete left-right transformation $P: \quad L \leftrightarrow R, \quad M \leftrightarrow M^\dagger$ which combined with the usual space inversion allows to build the parity transformation on the fields. If we ignore the transformations of eq. (1.1), the largest possible global symmetry of the low-energy theory is given by the requirement of maintaining for the transformed variables L', R' and M' the character of $SU(2)$ elements, or $G_{max} = [SU(2) \otimes SU(2)]^3$, consisting of three independent $SU(2) \otimes SU(2)$ factors, acting on each of the three variables separately. As we shall see, it happens that, for specific choices of the parameters of the theory, the symmetry G' gets enlarged to G_{max} [6].

The most general $G' \otimes P$ invariant lagrangian is given by [7]

$$L_G = -\frac{v^2}{4}[a_1 I_1 + a_2 I_2 + a_3 I_3 + a_4 I_4] \tag{1.2}$$

plus the kinetic terms L_{kin}. The four invariant terms I_i $(i = 1, ...4)$ are given by:

$$I_1 = tr[(V_0 - V_1 - V_2)^2] \quad I_2 = tr[(V_0 + V_2)^2] \quad I_3 = tr[(V_0 - V_2)^2] \quad I_4 = tr[V_1^2] \tag{1.3}$$

where

$$V_0^\mu = L^\dagger D^\mu L \quad V_1^\mu = M^\dagger D^\mu M \quad V_2^\mu = M^\dagger (R^\dagger D^\mu R) M \tag{1.4}$$

and the covariant derivatives are

$$D_\mu L = \partial_\mu L - L\mathbf{L}_\mu \quad D_\mu R = \partial_\mu R - R\mathbf{R}_\mu \tag{1.5}$$

$$D_\mu M = \partial_\mu M - M\mathbf{L}_\mu + \mathbf{R}_\mu M \tag{1.6}$$

where $\mathbf{L}_\mu(\mathbf{R}_\mu)$ are gauge fields of $SU(2)_{L(R)} \subset H'$ (instead of working with vector and axial-vector we work with these left and right combinations).

The kinetic terms are given by

$$L_{kin} = \frac{1}{g''^2} tr[F_{\mu\nu}(\mathbf{L})]^2 + \frac{1}{g''^2} tr[F_{\mu\nu}(\mathbf{R})]^2 \tag{1.7}$$

where g'' is the gauge coupling constant for the gauge fields $\mathbf{L}_\mu$ and $\mathbf{R}_\mu$, and $F_{\mu\nu}(\mathbf{L})$, $F_{\mu\nu}(\mathbf{R})$ are the usual field tensors.

The model we will consider is characterized by the following choice of parameters $a_4 = 0$, $a_2 = a_3$ [2, 6]. In order to discuss the symmetry properties that make such a choice natural it is useful to observe that the invariant I_1 could be re-written as $I_1 = -tr(\partial_\mu U^\dagger \partial^\mu U)$ with $U = LM^\dagger R^\dagger$ and the lagrangian as

$$L_G = \frac{v^2}{4}\{a_1 \ tr(\partial_\mu U^\dagger \partial^\mu U) + 2 \ a_2 \ [tr(D_\mu L^\dagger D^\mu L) + tr(D_\mu R^\dagger D^\mu R)]\} \tag{1.8}$$

Each of the three terms in the above expression is invariant under an independent $SU(2) \otimes SU(2)$ group

$$U' = \omega_L U \omega_R^\dagger, \quad L' = g_L L h_L, \quad R' = g_R R h_R \tag{1.9}$$

The overall symmetry is $G_{max} = [SU(2) \otimes SU(2)]^3$, with a part H' realized as a gauge symmetry. With the particular choice $a_4 = 0$, $a_3 = a_2$, as we see from eq. (1.8), the mixing between $\mathbf{L}_\mu$ and $\mathbf{R}_\mu$ is vanishing, and the new states are degenerate in mass. Moreover, as it follows from eq. (1.8), the longitudinal modes of the fields are entirely provided by the would-be Goldstone bosons in L and R. This means that the pseudoscalar particles remaining as physical states in the low-energy spectrum are those associated to U. They in turn can provide the longitudinal components to the W and Z particles, in an effective description of the electroweak breaking sector.

The peculiar feature of the model is that the new bosons are not coupled to those Goldstone bosons which are absorbed to give mass to W and Z. As a consequence the channels $W_L Z_L$ and $W_L W_L$ are not strongly enhanced as it usually happens in models with a strongly interacting symmetry breaking sector [3]. This also implies larger branching ratios of the new resonances into fermion pairs.

The coupling of the model to the electroweak $SU(2)_W \otimes U(1)_Y$ gauge fields is obtained via the minimal substitution

$$D_\mu L \to D_\mu L + \mathbf{W}_\mu L \quad D_\mu R \to D_\mu R + \mathbf{Y}_\mu R \quad D_\mu M \to D_\mu M \tag{1.10}$$

where

$$\begin{aligned} \mathbf{W}_\mu &= ig \tilde{W}_\mu^a \frac{\tau^a}{2} \quad \mathbf{Y}_\mu = ig' \tilde{Y}_\mu \frac{\tau^3}{2} \\ \mathbf{L}_\mu &= i \frac{g''}{\sqrt{2}} \tilde{L}_\mu^a \frac{\tau^a}{2} \quad \mathbf{R}_\mu = i \frac{g''}{\sqrt{2}} \tilde{R}_\mu^a \frac{\tau^a}{2} \end{aligned} \tag{1.11}$$

with g, g' the $SU(2)_W \otimes U(1)_Y$ gauge coupling constant and τ^a the Pauli matrices. We have used tilded quantities to reserve untilded variables for mass eigenstates.

By introducing the canonical kinetic terms for $\tilde{W}_\mu^a$ and $\tilde{Y}_\mu$ and going into the unitary gauge we get

$$\begin{aligned} \mathcal{L} &= -\frac{v^2}{4} \left[a_1 tr(\tilde{W}_\mu - \tilde{Y}_\mu)^2 + 2a_2 tr(\tilde{W}_\mu - \tilde{L}_\mu)^2 + 2a_2 tr(\tilde{Y}_\mu - \tilde{R}_\mu)^2 \right] \\ &+ \mathcal{L}^{kin}(\tilde{W}, \tilde{Y}, \tilde{L}, \tilde{R}) \end{aligned} \tag{1.12}$$

The standard model (SM) relations are obtained in the limit $g'' \gg g, g'$. Actually, for very large g'', the kinetic terms for the fields $\tilde{L}_\mu$ and $\tilde{R}_\mu$ drop out, and $\mathcal{L}$ reduces to the first term in eq. (1.12). This term reproduces precisely the mass term for the ordinary gauge vector bosons in the SM, provided we assume $a_1 = 1$. Finally let us consider the couplings to the fermions:

$$
\begin{aligned}
\mathcal{L}_{fermion} &= \overline{\psi}_L i\gamma^\mu\Big(\partial_\mu + ig\tilde{W}^a_\mu\frac{\tau^a}{2} + \frac{i}{2}g'(B-L)\tilde{Y}_\mu\Big)\psi_L \\
&+ \overline{\psi}_R i\gamma^\mu\Big(\partial_\mu + ig'\tilde{Y}_\mu\frac{\tau^3}{2} + \frac{i}{2}g'(B-L)\tilde{Y}_\mu\Big)\psi_R \qquad (1.13)
\end{aligned}
$$

where $B(L)$ is the baryon (lepton) number, and $\psi = (\psi_u, \psi_d)$. We have not introduced direct couplings to $\tilde{L}$ and $\tilde{R}$, so the new gauge bosons will couple to fermions only via mixing.

By separating the charged and the neutral gauge bosons, the quadratic lagrangian is given by:

$$
\begin{aligned}
\mathcal{L}^{(2)} &= \frac{v^2}{4}[(1+2a_2)g^2\tilde{W}^+_\mu\tilde{W}^{\mu-} + a_2 g''^2(\tilde{L}^+_\mu\tilde{L}^{\mu-} + \tilde{R}^+_\mu\tilde{R}^{\mu-}) \\
&- \sqrt{2}a_2 g g''(\tilde{W}^+_\mu\tilde{L}^{\mu-} + \tilde{W}^-_\mu\tilde{L}^{\mu+})] \\
&+ \frac{v^2}{8}[(1+2a_2)(g^2\tilde{W}_3^2 + g'^2\tilde{Y}^2) + a_2 g''^2(\tilde{L}_3^2 + \tilde{R}_3^2) \\
&- 2gg'\tilde{W}_{3\mu}\tilde{Y}^\mu - 2\sqrt{2}a_2 g''(g\tilde{W}_3\tilde{L}^\mu_3 + g'\tilde{Y}_\mu\tilde{R}^\mu_3)] \qquad (1.14)
\end{aligned}
$$

Therefore the $R^\pm$ fields are unmixed and their mass can be easily read: $M_{R^\pm} \equiv M = vg''\sqrt{a_2}/2$. All the other heavy fields have degenerate mass M in the large g'' limit (for this reason we call this model Degenerate BESS), and W and Z masses get corrections of order $(g/g'')^2$ [2]. We will parameterize the model by using, in addition to the SM parameters, M and g/g''.

By using eq. (1.14) one can show that, at the leading order in q^2/M^2, the contribution of the model to all ϵ parameters [8] is equal to zero [2]. This is due to the fact that in the $M \to \infty$ limit, this model decouples. We can perform the low-energy limit at the next-to-leading order and study the virtual effects of the heavy particles. Working at the first order in $1/g''^2$ we get $\epsilon_1 = -(c_\theta^4+s_\theta^4)/(c_\theta^2)\, X$, $\epsilon_2 = -c_\theta^2\, X$, $\epsilon_3 = -X$ with $X = 2(M_Z^2/M^2)(g/g'')^2$. All these deviations are of order X which contains a double suppression factor M_Z^2/M^2 and $(g/g'')^2$. The sum of the SM contributions, functions of the top and Higgs masses, and of these deviations has to be compared with the experimental values for the ϵ parameters, determined from the all available LEP data and the M_W measurement at Tevatron [9]: $\epsilon_1 = (3.8 \pm 1.5)\cdot 10^{-3}$, $\epsilon_2 = (-6.4 \pm 4.2)\cdot 10^{-3}$, $\epsilon_3 = (4.6 \pm 1.5)\cdot 10^{-3}$. Taking into account the SM values $(\epsilon_1)_{SM} = 4.4\cdot 10^{-3}$, $(\epsilon_2)_{SM} = -7.1\cdot 10^{-3}$, $(\epsilon_3)_{SM} = 6.5\cdot 10^{-3}$ for $m_{top} = 180\ GeV$ and $m_H = 1000\ GeV$, we find, from the combinations of the previous experimental results, the 90% C.L. limit on g/g'' versus the mass M given in Fig. 1. The allowed region is the one below the solid line.

2 e^+e^- future colliders

In this section we will discuss the sensitivity of the model at LEP2 and future e^+e^- linear colliders, for different options of total centre of mass energies and luminosities.

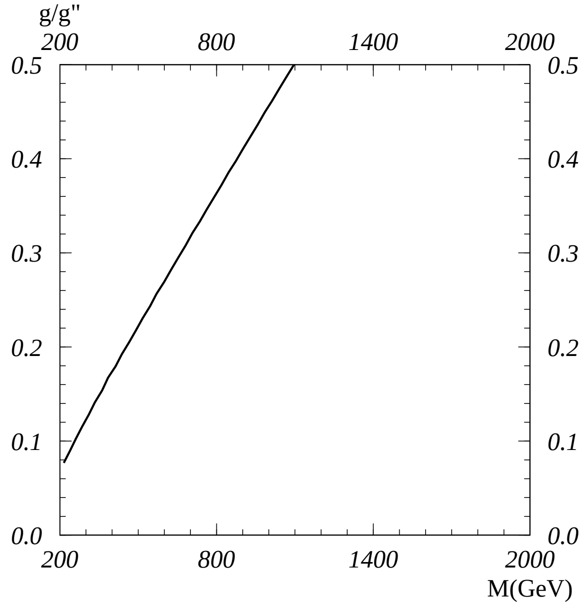

Fig. 1 - *90% C.L. contour on the plane (M, g/g'') obtained by comparing the values of the ϵ parameters from the model to the experimental data from LEP. The allowed region is below the curve.*

Cross-sections and asymmetries for the channel $e^+e^- \to f^+f^-$ and $e^+e^- \to W^+W^-$ in the SM and in the degenerate BESS model at tree level have been studied [2]. The BESS states relevant for the analysis at e^+e^- colliders are L_3 and R_3. Their coupling to fermions can be found in [2]. We will not consider the direct production of R_3 and L_3 from e^+e^-, but rather their indirect effects in the $e^+e^- \to f^+f^-$ and $e^+e^- \to W^+W^-$ cross-sections. In the fermion channel the study is based on the following observables: the total hadronic ($\mu^+\mu^-$) cross-sections σ^h (σ^μ), the forward-backward and left-right asymmetries $A_{FB}^{e^+e^-\to\mu^+\mu^-}$, $A_{FB}^{e^+e^-\to\bar{b}b}$, $A_{LR}^{e^+e^-\to\mu^+\mu^-}$, $A_{LR}^{e^+e^-\to h}$ and $A_{LR}^{e^+e^-\to\bar{b}b}$. At LEP2 we can add to the previous observables the W mass measurement. The result of this analysis shows that LEP2 will not improve considerably the existing limits [10].

In Fig. 2 we present the 90% C.L. contour on the plane (M, g/g'') from e^+e^- at $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80 fb^{-1}$ for various observables. The dashed-dotted line represents the limit from σ^h with an assumed relative error of 2%; the dashed line near to the preceeding one is σ^μ (relative error 1.3%), the dotted line is A_{FB}^μ (error 0.5%) and the uppermost dashed line is A_{FB}^b (error 0.9%).

As it is evident more stringent bounds come from the cross-section measurements. Asymmetries give less restrictive bounds due to a compensation between the L_3 and R_3 exchange. By combining all the deviations in the previously considered observables we get the limit shown by the continuous line.

Polarized electron beams allow to get further limit in the parameter space as shown in Fig. 3. We neglect the error on the measurement of the polarization and use a polarization value equal to 0.5. The dashed-dotted line represents the limit from A_{LR}^μ (error 0.6%), the dashed line from A_{LR}^h (error 0.4%), the dotted line from A_{LR}^b (error

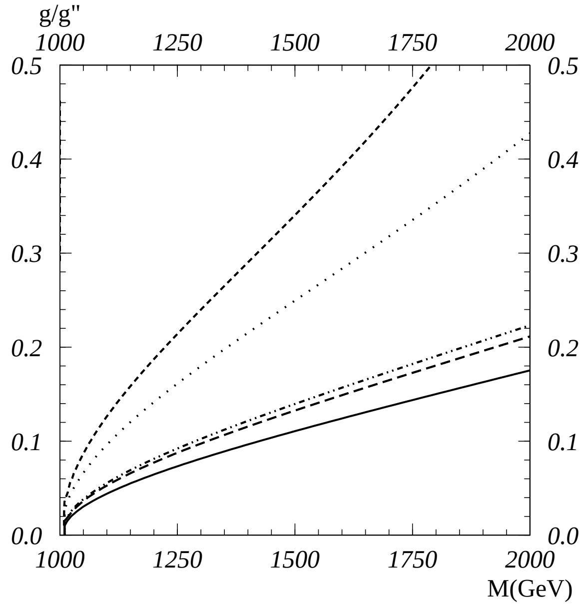

Fig. 2 - *90% C.L. contour on the plane $(M,\ g/g'')$ from e^+e^- at $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80 fb^{-1}$ from unpolarized observables. Allowed regions are below the curves. (Dashed-dotted is σ^h, dashed is σ^μ, dotted is A^μ_{FB}, the uppermost dashed is A^b_{FB}, continuous is for all combined).*

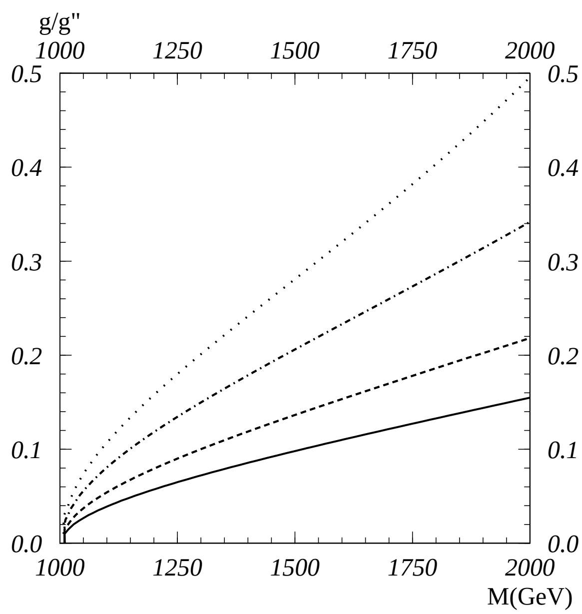

Fig. 3 - *90% C.L. contour on the plane $(M,\ g/g'')$ from e^+e^- at $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80 fb^{-1}$ from polarized observables. Allowed regions are below the curves. (Dashed-dotted is A^μ_{LR}, dashed is A^h_{LR}, dotted is A^b_{LR}, continuous is for all unpolarized and polarized combined).*

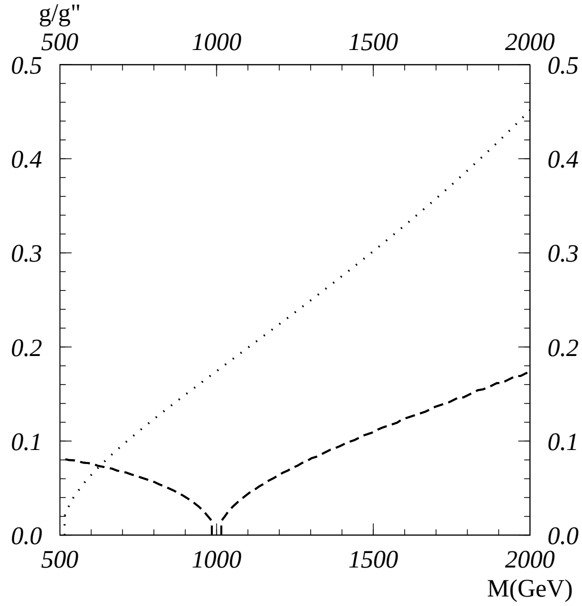

Fig. 4 - *90% C.L. contour on the plane* $(M,\ g/g'')$ *from* e^+e^- *at* $\sqrt{s} = 500\ GeV$ *with an integrated luminosity of* $20 fb^{-1}$ *and* $\sqrt{s} = 1000\ GeV$ *with an integrated luminosity of* $80 fb^{-1}$. *Allowed regions are below the curves.*

1.1%). Combining all the polarized and unpolarized beam observables we get the bound shown by the continuous line. In conclusion a substantial improvement with respect to the LEP bounds, even without polarized beams is obtained.

In Fig. 4 a combined picture of the 90% C.L. contours on the plane $(M,\ g/g'')$ from e^+e^- at two values of $\sqrt{s}$ is shown. The dotted line represents the limit from the combined unpolarized observables at $\sqrt{s} = 500\ GeV$ with an integrated luminosity of $20 fb^{-1}$; the dashed line is the limit from the combined unpolarized observables at $\sqrt{s} = 1000\ GeV$ with an integrated luminosity of $80 fb^{-1}$. As expected increasing the energy of the collider and rescaling the integrated luminosity result in stronger bounds on the parameter space.

The WW final state, considering the observables given in [2] has been also studied. However the new channel does not modify the strong limits obtained using the fermion final state. This is because the degenerate model has no strong enhancement of the WW channel, present in the usual strong electroweak models. For example, this is the most important channel for the BESS model with only vector resonances [11].

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